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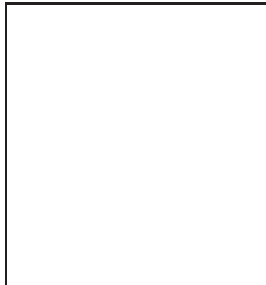
THE COLOR DIPOLE APPROACH TO THE DRELL-YAN PROCESS IN PA COLLISIONS

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In the target rest frame and at high energies, Drell-Yan (DY) dilepton production looks like bremsstrahlung of massive photons, rather than parton annihilation. The projectile quark is decomposed into a series of Fock states. Configurations with fixed transverse separations are interaction eigenstates for pp scattering. The DY cross section can then be expressed in terms of the same color dipole cross section as DIS. This approach is especially suitable to describe nuclear effects, since it allows to apply Glauber multiple scattering theory. We go beyond the Glauber eikonal approximation by taking into account transitions between interaction eigenstates. We calculate nuclear shadowing at large Feynman- x_F for DY in proton-nucleus collisions, compare to existing data from E772 and make predictions for RHIC. Nuclear effects on the transverse momentum distribution are also investigated.

1 Introduction

Although cross sections are Lorentz invariant, the partonic interpretation of the microscopic process depends on the reference frame. As pointed out by one of the authors, in the target rest frame DY dilepton production should be treated as bremsstrahlung, rather than parton annihilation¹ (see also²). The space-time picture of the DY process in the target rest frame is illustrated in fig. 1. A quark (or an antiquark) from the projectile hadron radiates a virtual photon on impact on the target. The radiation can occur before or after the quark scatters off the target. Only the latter case is shown in fig. 1.

A salient feature of the rest frame picture of DY dilepton production is that at high energies and in impact parameter space the DY cross section can be formulated in terms of the same dipole cross section as low- x_{Bj} DIS.

The color dipole approach to the DY process provides a convenient alternative to the well known parton model, in particular, it is especially appropriate to describe nuclear effects^{1,3}.

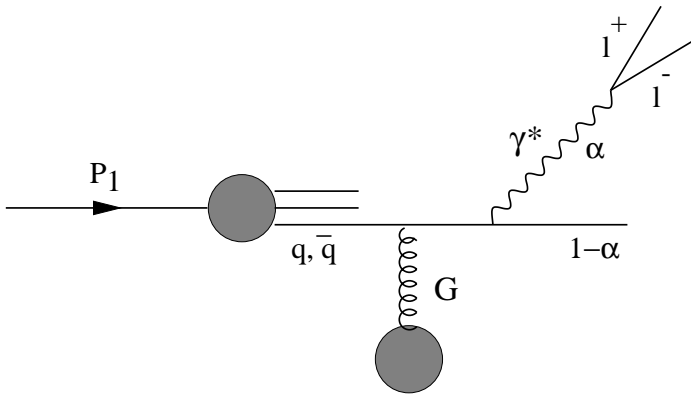


Figure 1: A quark (or an antiquark) inside the projectile hadron scatters off the target color field and radiates a massive photon, which subsequently decays into the lepton pair. The photon can also be radiated before the quark hits the target. Here, α is the longitudinal momentum fraction of the quark carried by the photon.

2 DY dilepton production in pp scattering

The cross section for radiation of a virtual photon from a quark after scattering on a proton, can be written in factorized light-cone form^{1,2,3},

$$\frac{d\sigma(qp \rightarrow \gamma^* X)}{d \ln \alpha} = \int d^2 \rho |\Psi_{\gamma^* q}^{T,L}(\alpha, \rho)|^2 \sigma_{q\bar{q}}(x_2, \alpha \rho), \quad (1)$$

similar to the case of DIS. Here, $\sigma_{q\bar{q}}$ is the cross section⁴ for scattering a $q\bar{q}$ -dipole off a proton which depends on the $q\bar{q}$ separation $\alpha\rho$, where ρ is the photon-quark transverse separation and α is the fraction of the light-cone momentum of the initial quark taken away by the photon. We use the standard notation for the kinematical variables, $x_1 - x_2 = x_F$, $\tau = M^2/s = x_1 x_2$, where x_F is the Feynman variable, s is the center of mass energy squared of the colliding protons and M is the dilepton mass. In (1) T stands for transverse and L for longitudinal photons.

The physical interpretation of (1) is similar to the DIS case. The projectile quark is expanded in the interaction eigenstates. We keep here only the first eigenstate,

$$|q\rangle = \sqrt{Z_2} |q_{bare}\rangle + \Psi_{\gamma^* q}^{T,L} |q\gamma^*\rangle + \dots, \quad (2)$$

where Z_2 is the wavefunction renormalization constant for fermions. In order to produce a new state the interaction must distinguish between the two Fock states, *i.e.* they have to interact differently. Since only the quarks interact in both Fock components the difference arises from their relative displacement in the transverse plane. If ρ is the transverse separation between the quark and the photon, the $\gamma^* q$ fluctuation has a center of gravity in the transverse plane which coincides with the impact parameter of the parent quark. The transverse separation between the photon and the center of gravity is $(1 - \alpha)\rho$ and the distance between the quark and the center of gravity is correspondingly $\alpha\rho$. Therefore, the argument of $\sigma_{q\bar{q}}$ is $\alpha\rho$. More discussion can be found in⁵.

The transverse momentum distribution of DY pairs can also be expressed in terms of the dipole cross section³. The differential cross section is given by the Fourier integral

$$\begin{aligned} \frac{d\sigma(qp \rightarrow \gamma^* X)}{d \ln \alpha d^2 q_\perp} &= \frac{1}{(2\pi)^2} \int d^2 \rho_1 d^2 \rho_2 \exp[i\vec{q}_\perp \cdot (\vec{\rho}_1 - \vec{\rho}_2)] \Psi_{\gamma^* q}^*(\alpha, \vec{\rho}_1) \Psi_{\gamma^* q}(\alpha, \vec{\rho}_2) \\ &\times \frac{1}{2} \{ \sigma_{q\bar{q}}(x_2, \alpha\rho_1) + \sigma_{q\bar{q}}(x_2, \alpha\rho_2) - \sigma_{q\bar{q}}(x_2, \alpha(\vec{\rho}_1 - \vec{\rho}_2)) \}. \end{aligned} \quad (3)$$

after integrating this expression over the transverse momentum q_\perp of the photon, one obviously recovers (1).

The LC wavefunctions can be calculated in perturbation theory and are well known^{2,5}. The dipole cross section on the other hand is largely unknown. Only at small distances ρ it can

be expressed in terms of the gluon density. However, several successful parameterizations exist in the literature, describing the entire function $\sigma_{q\bar{q}}(x, \rho)$, without explicitly taking into account the QCD evolution of the gluon density. We use the parameterization by Golec-Biernat and Wüsthoff⁶ for our calculations, fig. 2. This parameterization vanishes $\propto \rho^2$ at small distances, as implied by color transparency⁴ and levels off exponentially at large separations. The data in fig. 2 are quite well described without any K -factor, which does not appear in this approach since higher order corrections are supposed to be parameterized in $\sigma_{q\bar{q}}(x_2, \rho)$.

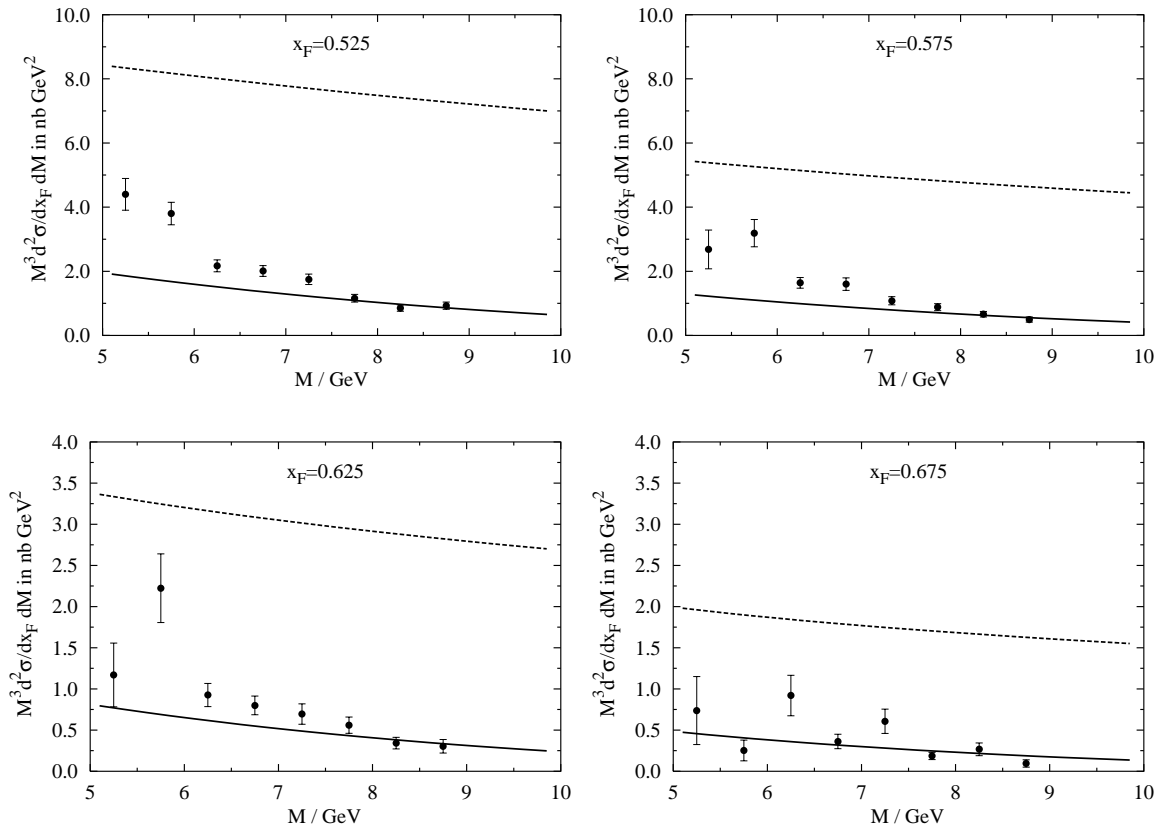


Figure 2: The points show the measured DY cross section in p^2H scattering from E772. The curves are calculated without any further fitting procedure. The solid curves are calculated at the same kinematics as the data points (center of mass energy $\sqrt{s} = 38.8$ GeV). The dashed curves are calculated for RHIC energy, $\sqrt{s} = 500$ GeV.

3 Proton-nucleus (pA) scattering

Shadowing in DY is an interference phenomenon due to multiple scattering of the projectile quark inside the nucleus. In the target rest frame, where DY dilepton production is bremsstrahlung of massive photons, shadowing is the Landau-Pomeranchuk-Migdal (LPM) effect. These interferences occur (fig. 1), because photons radiated at different longitudinal coordinates z_1 and z_2 are not independent of each other. Thus, the amplitudes have to be added coherently. Destructive interferences can occur only if the longitudinal distance $z_2 - z_1$ is smaller than the so called coherence length l_c , which is the time needed to distinguish between a quark and a quark with a γ^* nearby. It is given by the uncertainty relation,

$$l_c = \frac{1}{\Delta P^-} = \frac{1}{m_N x_2} \frac{(1 - \alpha)M^2}{q_{\perp}^2 + (1 - \alpha)M^2 + \alpha^2 m_q^2}. \quad (4)$$

Here, ΔP^- is the light-cone energy denominator for the transition $q \rightarrow q\gamma^*$ and q_\perp is the relative transverse momentum of the γ^*q Fock state. For $z_1 - z_2 > l_c$, the radiations are independent of each other.

An immediate consequence of this is that l_c has to be larger than the mean distance between two scattering centers in the nucleus (~ 2 fm in the nuclear rest frame). Otherwise, the projectile quark could not scatter twice within the coherence length and no shadowing would be observed.

We develop a Green function technique³, which allows one to resum all multiple scattering terms, similar to Glauber theory, and in addition treats the coherence length exactly. The formalism is equivalent to the one proposed in⁸ for the LPM effect in QED. Our general expression for the nuclear DY cross section reads

$$\begin{aligned} \frac{d\sigma(qA \rightarrow \gamma^*X)}{d\ln\alpha} &= A \frac{d\sigma(qp \rightarrow \gamma^*X)}{d\ln\alpha} - \frac{1}{2} \text{Re} \int d^2b \int_{-\infty}^{\infty} dz_1 \int_{z_1}^{\infty} dz_2 \int d^2\rho_1 \int d^2\rho_2 \\ &\times \left[\Psi_{\gamma^*q}(\alpha, \rho_2) \right]^* \rho_A(b, z_2) \sigma_{q\bar{q}}(x_2, \alpha\rho_2) G(\vec{\rho}_2, z_2 | \vec{\rho}_1, z_1) \\ &\times \rho_A(b, z_1) \sigma_{q\bar{q}}(x_2, \alpha\rho_1) \Psi_{\gamma^*q}(\alpha, \rho_1). \end{aligned} \quad (5)$$

The first term is just A times the single scattering cross section, where A is the nuclear mass number. The second term is the shadowing correction. The impact parameter is b and the nuclear density is ρ_A . The Green function G describes, how the bremsstrahlungs-amplitude at z_1 interferes with the amplitude at z_2 .

To make the meaning of Eq. 5 more clear, let us first consider a limiting case for G . In the simplest case, the coherence length, Eq. 4, is infinitely long and only the double scattering term is taken into account. Then $G(\vec{\rho}_2, z_2 | \vec{\rho}_1, z_1) = \delta^{(2)}(\vec{\rho}_1 - \vec{\rho}_2)$ and one of the ρ integrations can be performed. The δ -function means that at very high energy (infinite coherence length) the transverse size of the γ^*q Fock-state does not vary during propagation through the nucleus, it is frozen due to Lorentz time dilatation. Furthermore, partonic configurations with fixed transverse separations in impact parameter space were identified a long time ago⁴ in QCD as interaction eigenstates. This is the reason, why we work in coordinate space. Namely, in coordinate space, all multiple scattering terms can be resummed and in the limit of infinite l_c one obtains

$$G^{\text{frozen}}(\vec{\rho}_2, z_2 | \vec{\rho}_1, z_1) = \delta^{(2)}(\vec{\rho}_1 - \vec{\rho}_2) \exp\left(-\frac{\sigma_{q\bar{q}}(x_2, \rho_1)}{2} \int_{z_1}^{z_2} dz \rho_A(b, z)\right). \quad (6)$$

The frozen approximation is identical to eikonalization of the dipole cross section in Eq. (1). Thus, the impact parameter representation allows a very simple generalization from a proton to a nuclear target, provided the coherence length is infinitely long.

At Fermilab fixed-target energies ($\sqrt{s} = 38.8$ GeV for E772), this last condition is not fulfilled and one has to take a finite l_c into account. The problem is however, that l_c , Eq. 4, depends on the relative transverse momentum q_\perp of the γ^*q -fluctuation which is the conjugate variable to the size ρ of this Fock-state and therefore completely undefined in ρ -representation. The quantum mechanically correct way to treat the q_\perp^2 in Eq. 4 is to represent it by a two-dimensional Laplacian Δ_\perp in ρ -space. The Green function which contains the correct, finite coherence length and resums all multiple scattering terms fulfills a two-dimensional Schrödinger equation with an imaginary potential,

$$\begin{aligned} \left[i \frac{\partial}{\partial z_2} + \frac{\Delta_\perp(\rho_2) - \eta^2}{2E_q\alpha(1-\alpha)} + \frac{i}{2} \rho_A(b, z_2) \sigma_{q\bar{q}}(x_2, \alpha\rho_2) \right] G(\vec{\rho}_2, z_2 | \vec{\rho}_1, z_1) \\ = i\delta(z_2 - z_1) \delta^{(2)}(\vec{\rho}_2 - \vec{\rho}_1). \end{aligned} \quad (7)$$

For details of the derivation, we refer to³.

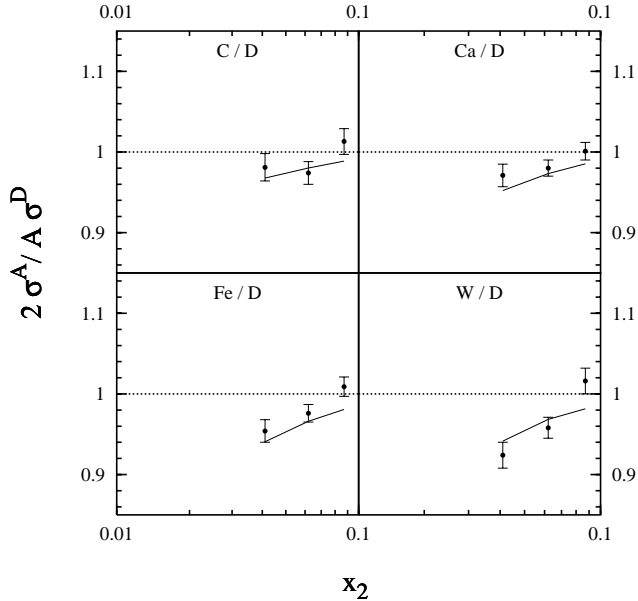


Figure 3: Comparison between calculations in the Green function technique and E772 data at center of mass energy $\sqrt{s} = 38.8$ GeV. for shadowing in DY. The calculations are performed at the mean values of the lepton pair mass. From the left to the right, these values are 5 GeV, 5.7 GeV, and 6.5 GeV.

The imaginary potential accounts for all higher order scattering terms. The Laplacian implies that the Green function is no longer proportional to a δ -function. This means the size of the γ^*q fluctuation is no longer constant during propagation through the nucleus. One can say that an eigenstate of size ρ_1 evolves to an eigenstate of size $\rho_2 \neq \rho_1$, so transitions between eigenstates occur.

Calculations¹⁰ with Eqs. 5 and 7 are compared to E772 data⁹ in fig. 3. Note that the coherence length l_c at E772 energy becomes smaller than the nuclear radius. Shadowing vanishes as x_2 approaches 0.1, because the coherence length becomes smaller than the mean internucleon separation. It is therefore important to have a correct description of a finite l_c in this energy range.

Nuclear effects on the q_\perp -differential cross section calculated at RHIC energy are shown in fig. 4. See³ for details of the calculation. The differential cross section is suppressed at small transverse momentum q_\perp of the dilepton, where large values of ρ dominate. This suppression vanishes at intermediate $q_\perp \sim 2$ GeV. In this region, one even observes an enhancement which reminds one of the Cronin effect. This enhancement is due to multiple scattering of the quark inside the nucleus. A nuclear target provides a larger momentum transfer than a proton target and harder fluctuations are freed, which leads to nuclear broadening. Note, that not the entire suppression at low q_\perp is due to shadowing. Some of the dileptons missing at low q_\perp reappear in this enhancement region. At very large transverse momentum nuclear effects vanish.

4 Summary

We express the DY cross section in terms of the cross section $\sigma_{q\bar{q}}$ for scattering a $q\bar{q}$ dipole off a proton. This is the same dipole cross section that appears in DIS. We can reasonably well describe low x_2 DY data from pp collisions⁷, without any free parameters and without a K factor.

At very high energy, the dipole approach is easily extended to nuclear targets by eikonalization. At lower fixed target energies (E772) the frozen approximation is no longer valid, because the size of a Fock state varies during propagation through the nucleus. Therefore, transitions between interaction eigenstates (i.e. partonic configurations with fixed transverse separations) occur.

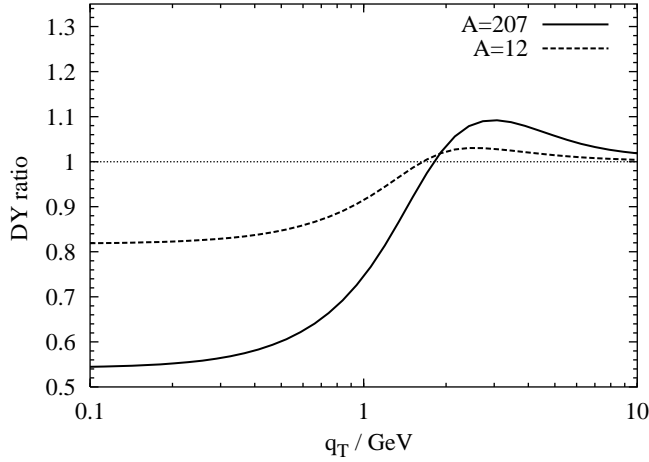


Figure 4: Nuclear effects on the DY transverse momentum distribution at RHIC energy ($\sqrt{s} = 200$ GeV) for carbon and lead. The figure shows the DY cross section ratio $(d\sigma^A/dx_F dM^2 d^2q_\perp)/(Ad\sigma^p/dx_F dM^2 d^2q_\perp)$ for dilepton mass $M = 5$ GeV and Feynman $x_F = 0.6$. Here, $q_\perp = q_T$ is the transverse momentum of the lepton pair.

We develop a Green function technique, which takes variations of the transverse size into account and resums all multiple scattering terms as well. Calculations with the Green function technique are in good agreement with DY shadowing data from E772. We have also calculated nuclear effects in the transverse momentum distribution of DY pairs at RHIC energy. The DY cross section is suppressed at low transverse momentum, but enhanced at intermediate $q_\perp \sim 2$ GeV. Nuclear effects vanish at very large q_\perp .

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